

# Transverse momentum spectra of hadrons produced in central heavy-ion collisions

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In-medium effects of transverse-mass distributions of quarks and gluons are considered assuming a possible local equilibrium for colorless quark objects like mesons and baryons created in central A+A collisions. It is shown that the average square of the transverse momentum for these partons grows and then saturates when the initial energy increases. Within the quark-gluon string model it leads to the energy dependence of hadron transverse mass spectra which is similar to that observed in heavy ion collisions. Comparison with other scenarios is given.

## 1. INTRODUCTION

Searching for a new physics in heavy-ion collisions at AGS, SPS and RHIC energies has led to intense theoretical and experimental activities in this field of research [1]. In this respect the search for signals of a possible transition of hadrons into the QCD predicted phase of deconfined quarks and gluons, quark-gluon plasma (QGP), is of particular interest. Among many signals of the QGP formation, one of the earliest is based on the relation of the thermodynamic variables, temperature and entropy, to the observable quantities, the average transverse momentum and multiplicity, respectively. Shuryak [2, 3] and van Hove [4] were first to put forth arguments that the transverse momentum as a function of multiplicity reflects general properties of the equation of state of hot and dense matter. In the case of the first order phase transition this relation is quite specific due to inevitable formation of the coexistence phase of quark-gluon plasma and hadrons, the mixed phase (MP). Experimental detection of the QGP phase and the MP in A+A collisions is a non-trivial task because of smallness of the space-time volume of the hot and dense system and possible contributions of hadronic processes simulating signals of QGP and MP formation. Nevertheless, recent experimental study of the transverse-mass spectra of kaons from central Au+Au and Pb+Pb collisions revealed "anomalous" dependence on the incident energy. The effective transverse temperature (the inverse slope-parameter of the transverse mass distribution at the mid-

rapidity) rather fast increases with incident energy in the AGS domain [5], then saturates at the SPS energies [6] and approaching the RHIC energy region [7] increases again. In agreement with expectations this saturation was assumed to be associated with the deconfinement phase transition and indication on the MP [8, 9]. This assumption on the first order phase transition was indirectly confirmed by the fact that two independent transport models, based on hadron-string dynamics, UrQMD and HSD, failed to reproduce the observed behavior of the effective kaon temperature [10, 11]. Though the calculated transverse mass spectra are in reasonable agreement with the experimental results for  $pp$ ,  $pA$  and light nuclei (central C+C and Si+Si) collisions, both models strongly underestimate the transverse energy  $T$  above energy  $\sim 5$  AGeV and are not able to describe the rapid increase of the inverse slope parameter for collisions in the AGS domain and its subsequent flattening in the low energy SPS regime. The failure was attributed to a lack of pressure. This additional pressure must be generated in the early nonhadronic phase of the collision because the strong hadronic interactions in the later stages do not produce it [10, 11]. The anomalous effective-temperature behavior was quite successfully reproduced within hydrodynamic model with equation of state involving the phase transition [12]. However, this result is not very convincing since to fit data the required incident-energy dependence of the freeze-out temperature should closely repeat the shape of the corresponding effective kaon temperature and thereby the problem of the observed anomalous inverse-slope dependence is readdressed to the problem of the freeze-out temperature.

In addition, as was shown in [13], these experimental data within the considered energy range up to  $\sim 160$  AGeV may be reasonably described in the 3-fluid relativistic hydrodynamic model with the pure hadronic equation of state. Note that this 3-fluid model also well describes a large variety of global observables [14]. This implies that a hydrodynamic collective motion in the expansion stage, which is absent in transport codes, plays an important role in the anomaly discussed. It is of interest that recently a step to taking into account some collective motion in the microscopic description has been done in [15] where a hadronic transport model has been extended to include three-body collisions. The agreement between the transport model and experiment for the kaon transverse-temperature excitation function was noticeably improved.

In this paper we propose another way to introduce a collectivity effect in nuclear system via some in-medium effect. We consider the temperature dependence of quark distribution

functions inside a colorless quark-antiquark or quark-diquark system (like meson or baryon,  $h$ ) created in central A+A collisions. A contribution of this effect into transverse momentum spectra of hadron is estimated and it is shown that it results in larger values of the inverse slope parameter and therefore in broadening of the transverse mass spectra.

## 2. QUARK IN A HADRON EMBEDDED IN EQUILIBRATED MATTER

Let us assume the local equilibrium in a fireball of hadrons whose distribution function can be presented in the following relativistic invariant form:

$$f_h^A = C_T \{1 \pm \exp((p_h \cdot u - \mu_h)/T)\}^{-1}, \quad (1)$$

where  $p_h$  is the four-momentum of the hadron, the four-velocity of the fireball in the proper system is  $u = (1, 0, 0, 0)$ , the sign "+" is for fermions and "-" is for bosons,  $\mu_h$  is the baryon chemical potential of the hadron  $h$ ,  $T$  is the local temperature, and  $C_T$  is the  $T$ -dependent normalization factor. The distribution function of constituent quarks inside  $h$  which is in local thermodynamic equilibrium with the surrounding nuclear matter,  $f_q^A(x, \mathbf{p}_t)$ , can be calculated using the procedure suggested for a free hadron in Ref.[16]. So, the quark distribution in a free hadron moving with the momentum  $P$  is given by the convolution

$$f_{q_v}^h(p_z, \mathbf{p}_t) = \int dp_{1z} d^2 p_{1t} q_v(p_z, \mathbf{p}_t) q_r(p_{1z}, \mathbf{p}_{1t}) \delta(p_z + p_{1z} - P) \delta^{(2)}(\mathbf{p}_{1t} + \mathbf{p}_t), \quad (2)$$

where  $q_v(p_z, \mathbf{p}_t)$  is related with the probability to find the valence quark with longitudinal momentum  $p_z$  and transverse momentum  $\mathbf{p}_t$  in the hadron, whereas  $q_r(p_{1z}, \mathbf{p}_{1t})$  is the probability that all the other hadron constituents (one or two valence quarks plus any number of quark-antiquark  $q\bar{q}$  pairs and gluons) carry the total longitudinal momentum  $p_{1z}$  and the total transverse momentum  $\mathbf{p}_{1t}$ .

Introducing the new variables  $x = p_z/P$ ,  $x_1 = p_{1z}/P$ , which at large  $P$  become the relativistic invariant light cone variables, we can rewrite eq.(2) (for example, for the valence quark distribution) in the following equivalent form [17, 18]:

$$f_{q_v}^h(x, p_t) = \int_0^1 dx_1 \int d^2 p_{1t} q_v(x, \mathbf{p}_t) q_r(x_1, \mathbf{p}_{1t}) \delta(x + x_1 - 1) \delta^{(2)}(\mathbf{p}_{1t} + \mathbf{p}_t), \quad (3)$$

Assuming the factorization hypothesis  $q_{v,r}(x, \mathbf{p}_t) = q_{v,r}(x) g_{v,r}(\mathbf{p}_t)$  and integrating eq.(3) with respect to  $dx_1$  and  $d^2 p_{1t}$  we have

$$f_{q_v}^h(x, p_t) = f_{q_v}^h(x) g_{q_v}^h(\mathbf{p}_t), \quad (4)$$

where

$$f_{qv}^h(x) = q_v(x) q_r(1-x) , \quad g_{qv}^h(\mathbf{p}_t) = g_v(\mathbf{p}_t) g_r(-\mathbf{p}_t) \quad (5)$$

For simplicity one can choose the normalized  $p_t$ -distribution  $g_q^h(\mathbf{p}_t)$  in the Gaussian form

$$g_q^h(\mathbf{p}_t) = \frac{\gamma_q}{\pi} \exp(-\gamma_q \mathbf{p}_t^2) . \quad (6)$$

Then we get

$$g_v(\mathbf{p}_t) = g_r(-\mathbf{p}_t) = \left(\frac{\gamma_q}{\pi}\right)^{1/2} \exp(-\frac{\gamma_q}{2} \mathbf{p}_t^2) . \quad (7)$$

Similarly to eq.(2), one can present the quark distribution inside the hadron  $h$  to be in local thermodynamic equilibrium with the surrounding matter,  $f_q^A(x, \mathbf{p}_t)$ , in the convolution form as

$$f_q^A(p_z, p_t) = \int dp_{1z} dp_{hz} d^2\mathbf{p}_{1t} d^2\mathbf{p}_{ht} \tilde{q}_v(p_z, \mathbf{p}_t) \tilde{q}_r(p_{1z}, \mathbf{p}_{1t}) f_h^A(p_{ht}, \mathbf{p}_{ht}) \times \delta(p_z + p_{1z} - p_{hz}) \delta^{(2)}(\mathbf{p}_t + \mathbf{p}_{1t} - \mathbf{p}_{ht}) . \quad (8)$$

In contrast to eq.(2) the distribution of the hadron  $h$  in a fireball  $f_h^A$  is included in eq.(8), therefore we integrate over the longitudinal and transverse momenta of  $h$ . In the general case, the functions  $q_v(p_{iz}, \mathbf{p}_{it})$  and  $q_r(p_{iz}, \mathbf{p}_{it})$  entering into eq.(2) may differ from  $\tilde{q}_v(p_{iz}, \mathbf{p}_{it})$  and  $\tilde{q}_r(p_{iz}, \mathbf{p}_{it})$  in eq.(8). Let us introduce the Feynman variables  $x = 2p_z^*/\sqrt{s'}$ ,  $x_1 = 2p_{1z}^*/\sqrt{s'}$ ,  $x_h = 2p_{hz}^*/\sqrt{s'}$ , where  $p_z^*, p_{1z}^*, p_{hz}^*$  are the longitudinal momenta and  $s'$  is some characteristic energy squared scale. Then eq.(8) can be reduced to the following form similar to eq.(3):

$$f_{qv}^A(x, p_t) = \int_0^1 dx_1 \int_0^1 dx_h \int d^2p_{1t} d^2p_{ht} \tilde{q}_v(x, \mathbf{p}_t) \tilde{q}_r(x_1, \mathbf{p}_{1t}) \times f_h^A(x_h, \mathbf{p}_{ht}) \delta(x + x_1 - x_h) \delta^{(2)}(\mathbf{p}_{1t} + \mathbf{p}_t - \mathbf{p}_{ht}) . \quad (9)$$

If the factorization hypothesis  $\tilde{q}_{v,r}(x, \mathbf{p}_t) = \tilde{q}_{v,r}(x) \tilde{g}_{v,r}(p_t)$  and the Gaussian form for  $\tilde{g}_{v,r}(\mathbf{p}_t)$  are assumed,

$$\tilde{g}_{x,r}(\mathbf{p}_t) = \left(\frac{\tilde{\gamma}_q}{\pi}\right)^{1/2} \exp(-\frac{\tilde{\gamma}_q}{2} \mathbf{p}_t^2) , \quad (10)$$

we can get the following expression for  $f_{qv}^A(x, p_t)$  (see **APPENDIX**) :

$$f_q^A(x, \mathbf{p}_t; T) = 2m_h T C_T \exp(-\frac{m_h - \mu_h}{T}) (1 + \frac{T}{m_h}) \frac{1}{\pi} \int_x^1 dx_h \tilde{q}_v(x) \tilde{q}_r(x_h - x) \times \tilde{\Gamma}_q(x_h) \exp(-\tilde{\Gamma}_q(x_h) p_t^2) , \quad (11)$$

where  $\tilde{m}_h(x_h) = \sqrt{m_h^2 + x_h s'/4}$ , by assumption  $\tilde{\gamma}_q = \gamma_q$  and

$$\tilde{\Gamma}_q(x_h) = \frac{\gamma_q(1 + \gamma_q \tilde{m}_h(x_h)T/2)}{1 + \gamma_q \tilde{m}_h(x_h)T} , \quad (12)$$

From the normalization relation for  $f_q^A(x, \mathbf{p}_t; T)$

$$\int_0^1 dx \int d^2 p_t f_q^A(x, \mathbf{p}_t; T) = 1 \quad (13)$$

we can find that  $C_T = (2m_h T(1 + T/m_h))^{-1} \exp[(m_h - \mu_h)/T]$ ,  $\tilde{q}_v(x) = q_v(x)$  and

$$q_r(1-x) = \int_x^1 dx_h \tilde{q}_r(x_h - x) \equiv \int_0^{1-x} \tilde{q}_r(y) dy . \quad (14)$$

Then eq.(11) reads

$$f_q^A(x, \mathbf{p}_t; T) = \frac{1}{\pi} \int_0^{1-x} dx_1 \tilde{q}_v(x) \tilde{q}_r(x_1) \tilde{\Gamma}_q(x_1 + x) \exp(-\tilde{\Gamma}_q(x_1 + x)p_t^2) . \quad (15)$$

As is evident from eq.(15), at vanishing temperature  $T = 0$  the quark distribution  $f_q^A$  reproduces the quark distribution in a free hadron.

Using the quark distribution  $f_{q_v}^A(x, p_t)$  we can calculate the average value for the transverse momentum squared for the quark  $\langle p_{q,t}^2 \rangle_h^A$  in  $h$  as a function of  $x$  and  $s_{NN}$ :

$$\langle p_{q,t}^2(x) \rangle_h^A = \frac{\int f_q^A(x, \mathbf{p}_t) p_t^2 d^2 p_t}{\int f_q^A(x, \mathbf{p}_t) d^2 p_t} . \quad (16)$$

Applying eq.(15), we have

$$\langle p_{q,t}^2(x) \rangle_h^A = C \int_0^{1-x} dx_1 \frac{\tilde{q}_r(x_1)(\langle p_t^2 \rangle_q^h + \tilde{m}_h(x_1 + x)T)}{1 + \tilde{m}_h(x_1 + x)T/(2 \langle p_t^2 \rangle_q^h)} , \quad (17)$$

where  $\langle p_t^2 \rangle_q^h = 1/\gamma_q$  is the average transverse momentum squared for a quark in the free hadron and

$$C^{-1} = \int_0^{1-x} dx_1 \tilde{q}_r(x_1) . \quad (18)$$

At  $x \simeq 0$  eq.(17) can be presented in the following equivalent form:

$$\langle p_{q,t}^2(x \simeq 0) \rangle_h^A = \tilde{C} \int_0^1 dx_1 \frac{\tilde{q}_r(x_1)(\langle p_t^2 \rangle_q^h + (\sqrt{m_h^2 + x_1^2 s'/4}T))}{1 + (\sqrt{m_h^2 + x_1^2 s'/4}T)/(2 \langle p_t^2 \rangle_q^h)} , \quad (19)$$

where

$$\tilde{C}^{-1} = \int_0^1 dx_1 \tilde{q}_r(x_1) . \quad (20)$$

In the general case the quark distribution in a nucleon  $f_q^N(x)$  at low momentum transfer when its  $Q^2$  QCD evolution can be neglected is presented as follows:

$$f_q^N(x) = C_q x^a (1-x)^b, \quad (21)$$

where  $C_q$  is the normalization factor. The parameters  $a$  and  $b$  can be extracted from the deep inelastic scattering or calculated within some quark models. Therefore, according to eq.(14), the function  $\tilde{q}_r(x_1)$  entering into eqs.(17,18) has the following form:

$$\tilde{q}_r(x_1) = b x_1^{b-1}. \quad (22)$$

Usually,  $b \geq 1.5$  (see below). So  $\tilde{q}_r(x_1)$  falls down very fast when  $x_1$  decreases from 1 to 0 and may be taken out of the integral in eq.(19) at  $x_1 = 1$  :

$$\langle p_{q,t}^2(x \simeq 0) \rangle_{h,appr.}^A \simeq \frac{\langle p_t^2 \rangle_q^h + T \sqrt{m_h^2 + s'/4}}{1 + T \sqrt{m_h^2 + s'/4} / (2 \langle p_t^2 \rangle_q^h)}, \quad (23)$$

As is evident from eq.(23) the quantity  $\langle p_{q,t}^2 \rangle_h^A$  depends on the energy  $\sqrt{s'}$  and temperature  $T$ . At any nonzero values of  $T$  it grows when  $\sqrt{s'}$  increases (because its derivative with respect to  $\sqrt{s'}$  is positive) and then saturates, reaching the asymptotic value about  $2 \langle p_t^2 \rangle_q^h$  at high energies. The integration over  $dx_1$  in eq.(19) does not change this result qualitatively.

To calculate  $\langle p_{q,t}^2(x) \rangle_h^A$  more accurately, we have to know the functions  $q_v(x) = \tilde{q}_v(x)$  and  $q_r(1-x)$  which is related to  $\tilde{q}_r(y)$ , see eqs.(14),(22). Parameters of these distributions can be defined by application of the quark-gluon string model (QGSM) [21, 22] based on the Regge asymptotic of quark distributions in a nucleon and  $1/N$  expansion in QCD [19, 20] ( $N$  is the number of flavors or colors). According to this model, the nucleon consists of a quark, diquark, and quark-antiquark see  $(q\bar{q})$ . For example, for the valence  $u$ -quark in the proton [22] we have

$$a = -\alpha_R(0) ; b = \alpha_R(0) - 2\alpha_N(0) \quad (24)$$

where  $\alpha_R(0) = 1/2$  is the reggeon intercept of the Regge trajectory,  $\alpha_N(0) = -0.5$  is the intercept of the nucleon Regge trajectory. Comparing eq.(21) and eq.(5) one can find that

$$u_v(x) = x^{-\alpha_R(0)} ; u_r(1-x) = (1-x)^{\alpha_R(0)-2\alpha_N(0)} \quad (25)$$

The similar form can be obtained for the  $x$ -distribution of the valence  $d$ -quark in the nucleon. Then, using eq.(19), the average transverse momentum squared can be estimated for  $u$ -quark

inside the proton which is in local equilibrium in the fireball  $\langle p_{u,t}^2 \rangle_p^A$ . This quantity at  $x \simeq 0$  is presented in Fig.1 as a function of  $\sqrt{s'}$ . The behavior of this quantity given by eq.(19) is similar to that discussed above in respect to the approximate expression (23).

### 3. TRANSVERSE MOMENTUM SPECTRA OF HADRONS FROM CENTRAL $A + A$ COLLISIONS

Let us now estimate the  $p_t$  distribution of the hadron  $h_1$  produced after collision of two hadrons one of them is locally equilibrated in a fireball. We shall explore the QGSM based on  $1/N$  expansion in QCD. Actually, this is the expansion of the QCD amplitude in different topologies [19, 20].<sup>1</sup> The first order term is the so-called planar graphs corresponding to the one-Reggeon exchange diagrams in the  $t$ -channel of the hadronic process. The second order term of this expansion is the so-called cylinder graphs related to the one-Pomeron exchange diagrams. The last ones make the main contribution to inclusive spectra of particles produced in inelastic hadronic processes. In Fig.2 the cylinder graphs for inelastic meson-nucleon (left diagram) and nucleon-nucleon (right diagram) inelastic processes are presented, see also Ref.[20]. According to this model, the colorless strings are formed between the antiquark/quark ( $\bar{q}/q$ ) in the colliding meson and the quark/diquark ( $q/qq$ ) in the colliding nucleon (left diagram of Fig.2), then, after their break,  $q\bar{q}$  pairs are created and fragmentate into a hadron  $h_1$ . The contribution of the cylinder graph (right diagram of Fig.2) to the inclusive spectrum  $\rho_{h_1}^A \equiv E_{h_1} \frac{d\sigma}{d^3p_{h_1}}$  of the hadron  $h_1$  from the collision of two nucleons can be presented as follows: [22, 26]:

$$\rho_{h_1}^A(x, \mathbf{p}_t; T) = \sigma_1 [F_q^{h_1}(x_+, p_t; T) F_{qq}^{h_1}(x_-, \mathbf{p}_t; T) / F_{qq}(0, p_t; T) + F_{qq}^{h_1}(x_+, p_t; T) F_q^{h_1}(x_-, \mathbf{p}_t; T) / F_q(0, p_t; T)] , \quad (26)$$

where  $\sigma_1$  is the cross section of the 2-chain production, corresponding to the  $s$ -channel discontinuity of the cylinder (one-Pomeron) graph. It is usually calculated within the quasi-

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<sup>1</sup> This model, the Quark-gluon String Model (QGSM) [21, 22] or the Dual Parton Model (DPM) [23] differs from the Lund String Model, see for example Ref.[24].

eikonal approximation [25];

$$F_{q(qq)}^{h_1}(x_{\pm}, \mathbf{p}_t, T) = \sum_{flavors} \int_{x_{\pm}}^1 dx_1 \int d^2 p_{1t} d^2 p_{2t} f_{q(qq)}^A(x_1, \mathbf{p}_{1t}; T) G_{q(qq)}^{h_1}\left(\frac{x_{\pm}}{x_1}, \mathbf{p}_{2t}\right) \times \delta^{(2)}(\mathbf{p}_{1t} + \mathbf{p}_{2t} - \mathbf{p}_t) . \quad (27)$$

Here  $x_{\pm} = \frac{1}{2}(\sqrt{x_t^2 + x^2} \pm x)$  and  $x_t = 2m_{h_1 t}/\sqrt{s_{hh}}$ , where  $s'$  has been associated with the energy squared  $s_{hh}$  of colliding hadrons,  $m_{h_1 t} = \sqrt{m_{h_1}^2 + p_{1t}^2}$ ;  $z^{-1}G_{q(qq)}^{h_1} = D_{q(qq)}^{h_1}$  is the fragmentation function (FF) of the quark  $q$  (diquark  $qq$ ) into the hadron  $h_1$ . Actually, the interaction function  $F_{q(qq)}^{h_1}(x_+, p_t, T)$  corresponds to the fragmentation of the upper quark/diquark ( see right diagram in Fig.2) into the hadron  $h_1$ , whereas  $F_q^{h_1}(x_-, p_t, T)$  corresponds to the fragmentation of the down diquark/quark into  $h_1$ . We can get a similar expression for the  $p_t$ -spectrum of the hadron  $h_1$  produced in the meson-nucleon collision by replacing the diquark  $qq$  with the antiquark  $\bar{q}$ . To calculate the transverse momentum spectrum of the hadron  $h_1$  in the central rapidity region, eq.(27), one needs to know the  $p_t$ -dependence of the fragmentation function  $D_q^{h_1}$ . We assume the same Gaussian dependence as eq.(6). However, the slope of this  $p_t$  dependence  $\gamma_c$  can differ from the slope  $\gamma_q$  for the constituent quark  $p_t$  distribution. For the mid rapidity region,  $x \simeq 0$ , eq.(26) can be rewritten in the following form:

$$\rho_{h_1}^A(x \simeq 0, \mathbf{p}_t; T) = \sigma_1[F_q^{h_1}(x \simeq 0, p_t; T)] + F_{qq}^{h_1}((x \simeq 0, p_t; T)) , \quad (28)$$

We have the following approximate equation for the averaged transverse momentum squared of the hadron  $h_1$  produced in the interaction of the nucleon with the nucleon locally equilibrated in a fireball:

$$\langle p_{h_1 t}^2 \rangle_{NN,appr.}^{AA} \simeq \frac{\tilde{\Gamma}_q(x \simeq 1) + \gamma_c}{\tilde{\Gamma}_q(x \simeq 1)\gamma_c} = \frac{1}{\tilde{\Gamma}_q(x \simeq 1)} + \frac{1}{\gamma_c} . \quad (29)$$

By substitution of  $\tilde{\Gamma}_q(x \simeq 1)$  given by eq.(12) at  $x_h = 1$  we finally have

$$\langle p_{h_1 t}^2 \rangle_{NN,appr.}^{AA} \simeq \frac{\langle p_t^2 \rangle_q^N + T\sqrt{m_N^2 + s_{hh}/4}}{1 + T\sqrt{m_N^2 + s_{hh}/4} / (2 \langle p_t^2 \rangle_q^N)} + \frac{\langle p_t^2 \rangle_q^N}{r} , \quad (30)$$

where  $r = \gamma_c/\gamma_q$ .

From comparison of eq.(30) and eq.(23) one can see that the energy dependence of the transverse momentum squared for the produced hadron,  $\langle p_{h_1 t}^2 \rangle_{NN,appr.}^{AA}$ , is qualitatively similar to that of the quark inside the hadron which is in local equilibrium in the fireball,



$\langle p_{q,t}^2 \rangle_h^A$ . At large values of  $\gamma_c \gg \gamma_q$ , the second term in eq.(30) can be neglected and  $\langle p_{h_1,t}^2 \rangle_{NN,appr}^{AA}$  as a function of  $\sqrt{s_{hh}}$  and  $T$  is close to  $\langle p_{q,t}^2 \rangle_{h,appr}^A$  given by eq.(23). In this case the hadron spectrum copies the quark spectrum. In fact,  $\gamma_c$  is larger than  $\gamma_q$  by factor  $r = 3 - 4$  as follows from inclusive  $p_t$ -spectra of hadrons produced in hadronic processes analyzed within the QGSM Ref.[26].

#### 4. RESULTS AND DISCUSSION

We estimated the average value of transverse momentum squared for  $K^+$ -mesons produced in nucleon-nucleon  $\langle p_{K^+,t}^2 \rangle_{NN}^{AA}$  and pion+nucleon  $\langle p_{K^+,t}^2 \rangle_{\pi N}^{AA}$  interactions of two hadrons one of them is thermodynamically locally equilibrated in a fireball created in the central  $A + A$  collision. This quantity was estimated as a function of  $\sqrt{s_{hh}}$  at  $T = 150$  MeV. for two cases when  $\gamma_c \gg \gamma_q$  and  $\gamma_c = 3\gamma_q$  [26] using eq.(29) for the inclusive spectrum at  $x \simeq 0$ . It is presented in Fig.3.

As is seen from Fig.3, the results obtained are sensitive to the mass value of the hadron which is in local equilibrium with the surrounding nuclear matter at  $\sqrt{s_{hh}} \leq 10$  GeV.

Collectivity due to thermal effects results in growth of  $\langle p_{K^+,t}^2 \rangle_{hh}^{AA}$  with energy for colliding hadrons in central A-A collisions similar to that for unbound quarks in hot matter,  $\langle p_{q,t}^2 \rangle_h^A$ . At larger energies  $\sqrt{s_{hh}}$  this quantity saturates. The saturation value for this quantity depends on the hadronization mechanism of quarks/diquarks to hadrons. A simple exponential estimate of  $p_t$ -spectra for produced hadrons  $h_1$  is used to parameterize experimental data:

$$\frac{dN}{dm_{h_1,t}^2 dy}|_{y=0} = C \exp(-\frac{m_{h_1,t}}{T^*}), \quad (31)$$

where the parameter  $T^*$  is extracted from fitting experimental data. There are data on the  $T^*$  values for different hadrons and different  $m_t$  domains: "low  $p_t$ " when  $m_{h_1,t} - m_{h_1} < 0.6$  GeV, and "high  $p_t$ ",  $0.6 < m_{h_1,t} - m_{h_1} < 1.6$  GeV (see for example Refs.[5-7]). At low  $p_t$  the  $m_t$  spectrum given by eq.(31) can be presented in the following approximate form:

$$\frac{dN}{dm_t^2 dy}|_{y=0} \simeq C \exp(-m_{h_1}) \exp(-\frac{p_{h_1,t}^2}{2m_{h_1} T^*}) \quad (32)$$

Actually, at small transverse momenta  $2m_{h_1} T^* \simeq \langle p_{h_1,t}^2 \rangle^{AA}$ , where  $\langle p_{h_1,t}^2 \rangle^{AA}$  is the transverse momentum squared for the hadron  $h_1$  produced in central  $A + A$  collisions. For

$K$ -mesons  $\langle p_{K,t}^2 \rangle^{AA} \simeq T^* GeV/c^2$ . The experimental data on the inverse slope  $T^*$  for  $K^+$ -mesons produced in central Au+Au (Pb+Pb) collisions as a function of the incident energy per nucleon show [5, 6] that  $T^*$  grows and saturates later on.

Our results presented in Fig.3 qualitatively demonstrate similar behavior for  $\langle p_{K^+,t}^2 \rangle_{hh}^{AA}$  as a function of  $\sqrt{s_{hh}}$ . One should emphasize that here  $\sqrt{s_{hh}}$  is the energy of a pair of colliding hadrons and it is not related directly to the initial energy of colliding heavy ions. These results are only an illustration of the collective effects assumed. In a real case, such binary interactions occur between various hadrons in a large range of temperatures.

As was noted in the introductional part, the broadening effect for  $m_t$ -spectra of the hadron produced in central A+A collisions observed at AGS, SPS and RHIC energies can qualitatively be explained using the assumption of possible creation of the QGP and a co-existing phase of quarks and hadrons [8, 9]. From available hydrodynamic calculations of the transverse inverse-slope excitation function [12, 13] it is not yet clear whether this behavior can be considered as a signal of the phase transition into the QGP. Microscopic transport models [10, 11] taking into account formation and decay of strings as well as the multiple rescatterings of hadrons are definitely not able to describe these data. An attempt to enhance the rescatterings was undertaken in Ref. [28] by including the Cronin effect. In the transport approach enhancement of the intrinsic quark transverse momentum spread  $\langle p_t^2 \rangle_q^A$  is simulated by increasing the average transverse momentum of quarks  $\langle p_t^2 \rangle_q$  with the number of previous collisions of primary nucleons  $N_{prev}$  as

$$\langle p_t^2 \rangle_q^A = \langle p_t^2 \rangle_q (1 + \alpha N_{prev}) , \quad (33)$$

where the parameter  $\alpha \approx 0.4$ . Now the description of spectra becomes rather good at the RHIC energies, improves essentially at the SPS energy of 160 AGeV, but does not show any significant change at 11 AGeV [28]. Consequently, the "pre-hadronic" Cronin effect, realized via eq.(33), is not responsible for the anomalous behavior of kaon slopes around AGS energies.

Another scenario of collectivity, color rope formation, was proposed in [29]. This color rope model assumes that in central A+A collisions several strings are produced, some of them on top of each other. The common chromo-electric field created by overlapping  $K$  single quark/antiquark sources may form a  $K$ -fold rope. The spread of hadron transverse mass distributions resulting from color string-rope decay is defined by the surface tension

parameter  $\kappa^K$  which is  $\sqrt{K}$  times larger than the appropriate parameter for the decay of a single string in  $N - N$  collisions,  $\kappa^K = \sqrt{K}\kappa$ . In a certain sense, this scenario is opposite to the proposed one: changing originates not from the initial "pre-hadronic" level, but rather comes from the final state as in-medium modification of the string break function. Similar effect for overlapping strings was estimated in Ref.[28]. Only a small increase in the inverse slope parameter at AGS energies was found because the string densities are low. At SPS and RHIC energies the model gives hardening of the spectra by about 15% [28].

## 5. CONCLUSION

We have found that the quark distribution in a hadron depends on the fireball temperature  $T$ . At any  $T$  the average transverse momentum squared of a quark grows and then saturates when  $\sqrt{s'}$  increases. Numerically this saturation property depends on  $T$ . The modification of initial quark distributions leads to a similar energy dependence for the average transverse momentum squared  $\langle p_{h_1,t}^2 \rangle_{hh}^{AA}$  of the hadron  $h_1$ . The saturation property for  $\langle p_{h_1,t}^2 \rangle_{hh}^{AA}$  depends also on the temperature  $T$  and it is very sensitive to the dynamics of hadronization. As an example, we estimated the energy dependence of the inverse slope of transverse mass spectrum of  $K$ -mesons produced in the interaction of two hadrons in the fireball created in central A+A collisions. It is qualitatively similar to the incident energy dependence of this quantity observed experimentally. We guess that our assumption on the thermodynamical equilibrium of hadrons given by eq.(1) can be applied for heavy nuclei only and not for the early interaction stage.

From the above discussion we see that the observed anomalous behavior of the kaon inverse slope in central A+A collisions is still puzzling. There are several scenarios which can be valid in various degrees, however, the final consistent solution of this puzzle is still absent. Its solution can partially be due to the proposed thermal mechanism of collectivity of hadrons. For final decision this and other effects should be taken into in a dynamical transport model. The solution of this "step-like" puzzle is an important point in the scientific programs on the future heavy-ion accelerators FAIR [30] and NICA [31]

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### Figure Captions

Fig.1. The energy dependence of the average transverse momentum squared for the  $u$ -quark in a proton in nuclear matter at temperature  $T$ .

Fig.2. The cylinder graph for the inelastic meson-nucleon processes (left) and the cylinder graph for the inelastic nucleon-nucleon reaction (right) [20]

Fig.3. The average transverse momentum squared of the  $K^+$ -meson produced from the interaction of two hadrons one of them is in the equilibrated fireball as a function of its energy  $\sqrt{s_{hh}}$  at  $T = 0.15$  GeV. Curves 1 and 2 correspond to  $\langle p_{K^+,t}^2 \rangle_{NN}^{AA}$  and  $\langle p_{K^+,t}^2 \rangle_{\pi N}^{AA}$  respectively when  $\gamma_c \gg \gamma_q$ , whereas curves 3 and 4 correspond to the same quantities when  $\gamma_c = 3\gamma_q$ . Line 5 corresponds to the average transverse momentum squared of  $K^+$  produced in free  $p + p$  collisions  $\langle p_t^2 \rangle_{K^+}^{NN} = 0.14 \text{ GeV}/c^2$ .

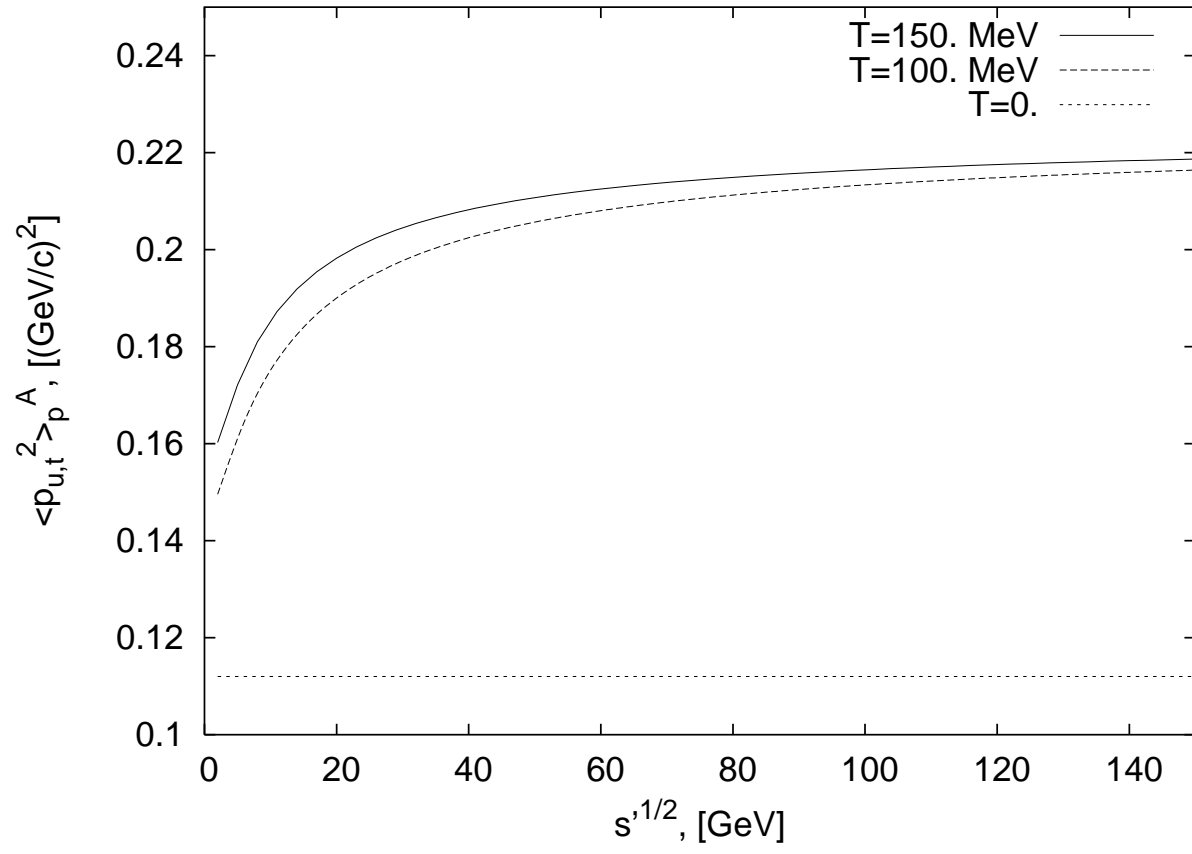


Figure 1.

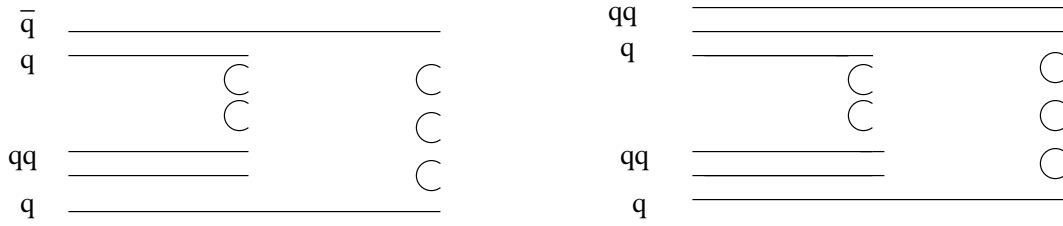


Figure 2.

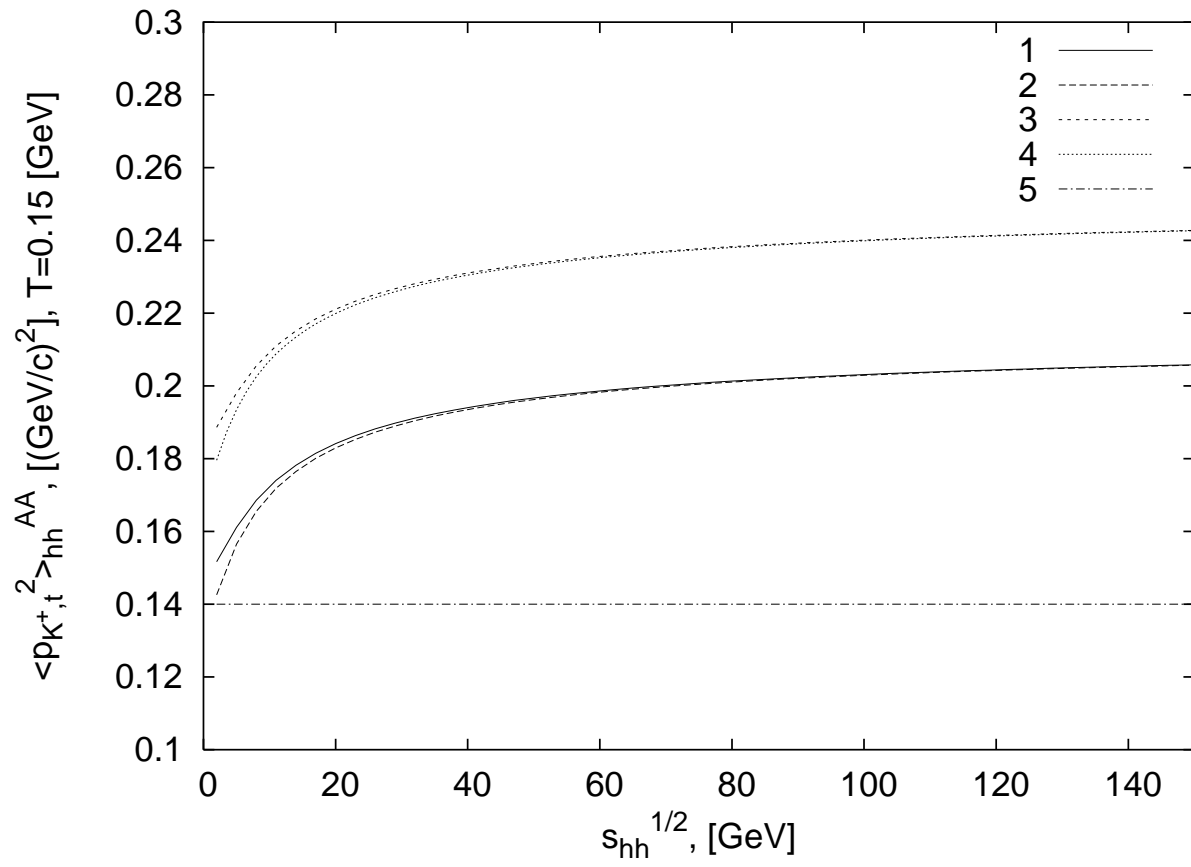


Figure 3.

## APPENDIX

Let us get eq.(11) using eqs.(9,7) and eq.(6). Eq.(9) can be also written in the following equivalent form

$$f_q^A(x, \mathbf{p}_t; T) = \frac{\gamma_q}{\pi} C_T \exp\left(\frac{\mu_h}{T}\right) \int_x^1 dx_h \tilde{q}_v(x) \tilde{q}_r(x_h - x) \exp(-\gamma'_q p_t^2) \times (34)$$

$$\times \frac{1}{(2\pi)^2} \int \exp(-\gamma'_q k_{1t}^2) \times \exp\left(-\frac{\sqrt{k_{2t}^2 + \tilde{m}_h^2(x_h)}}{T}\right) \exp(i\vec{b} \cdot (\vec{k}_{1t} + \vec{k}_{2t} - \vec{p}_t)) d^2 k_{1t} d^2 k_{2t} d^2 b ,$$

where  $\gamma'_q = \tilde{\gamma}_q/2$ ,  $\tilde{m}_h^2(x_h) = m_h^2 + x_h^2 s'/4$ . The first integral in eq.(34) reads

$$J_1 = \int \exp\left(-\frac{\sqrt{k_{2t}^2 + \tilde{m}_h^2(x_h)}}{T}\right) \exp(i\vec{b} \cdot \vec{k}_{2t}) d^2 k_{2t} \equiv 2\pi \int_0^\infty \exp\left(-\frac{\sqrt{k_{2t}^2 + \tilde{m}_h^2(x_h)}}{T}\right) \times (35)$$

$$\times J_0(bk_{2t}) k_{2t} dk_{2t} = 2\pi \frac{a_T}{(a_T^2 + b^2)^{3/2}} (1 + \tilde{m}_h(x_h) \sqrt{a_T^2 + b^2}) \exp(-m_h \sqrt{a_T^2 + b^2}) ,$$

where  $a_T = 1/T$ ,  $J_0(bk_{2t})$  is the Bessel function of order 0 depending on  $bk_{2t}$ .

Eq.(35) in the central region when  $b^2 < 1/T^2$  can be presented in the following approximate form:

$$J_1 = 2\pi T (\tilde{m}_h(x_h) + T) \exp(-b^2 \tilde{m}_h(x_h) T/2) . \quad (36)$$

Using now the form for  $J_1$  given by eq.(36) we can calculate the second integral in eq.(34)

$$\int \exp\left(-\frac{b^2}{4\gamma'_q}\right) \exp\left(-\frac{b^2 \tilde{m}_h T}{2}\right) \exp(-i\vec{b} \cdot \vec{p}_t) d^2 b = \frac{4\pi\gamma'_q}{(1 + 2\gamma'_q \tilde{m}_h T)} \exp\left(-\frac{\gamma'_q p_t^2}{1 + 2\gamma'_q \tilde{m}_h T}\right) . \quad (37)$$

Including all the terms staying in front of eq.(35) we get the form for  $f_q^A(x, \mathbf{p}_t; T)$  given by eq.(11). Note that by getting eq.(11) the term

$$(\tilde{m}_h(x_h) + T) \exp\left(-\frac{\tilde{m}_h(x_h)}{T}\right)$$

was moved out the integral in eq.(34) at  $x_h \simeq 0$  because the exponential function falls down very fast when  $x_h$  increases from 0 up to 1.